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Classical Nambu-Goldstone fields*

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Abstract

It is shown that a true Nambu-Goldstone (NG) boson develops a coherent long-range field whenever the charge associated with it that is carried by the other particles is not conserved in a macroscopic scale. The source of a NG field is the time rate of quantum number violation. If the lepton numbers are spontaneously broken at a scale below 1 TeV, the neutrino oscillation processes generate long-range majoron fields that are strong enough in Supernovae to modify the neutrino flavor dynamics. Two examples are given: NG fields may improve the adiabaticity of $\nu_e \leftrightarrow \nu_X$ transitions or cause *resonant* antineutrino oscillations otherwise impossible with solely weak interactions.

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1 Basic Features

The Goldstone theorem asserts that in a theory with spontaneous breaking of a global symmetry group some massless scalar bosons should exist – the so-called Nambu-Goldstone bosons – one per broken global symmetry. Good candidates for spontaneously broken quantum numbers are the partial lepton numbers L_e, L_μ, L_τ and total lepton number L. Indeed, they are conserved by the Standard Model (SM) interactions, but several experiments going now from solar to atmospheric neutrinos and laboratory oscillation experiments [1, 2] indicate that the lepton flavors are not conserved.

The NG bosons not only have zero mass but also no scalar potential terms such as ϕ^4. They interact with the fermion particles however, do not mediate long range forces. The reason lies in the very global symmetry which, broken by the vacuum still operates at the Lagrangian level, realized as a translation of the NG field, $\phi \to \phi + \alpha\,,\ \alpha = \text{constant}$. For that reason they only have derivative couplings [3] *i.e.*, the Lagrangian only depends on the derivatives $\partial_\mu \phi$. The ϕ equation of motion identifies with the conservation law associated to the $\phi \to \phi + \alpha$ symmetry,

$$\partial_\mu \partial^\mu \phi = -\partial_\mu J^\mu_\Lambda / V_\Lambda\,, \tag{1}$$

where V_Λ is essentially the scale of symmetry breaking and the current J^μ_Λ is determined in leading order by the quantum numbers of the other particles. For instance, if Λ is one of the lepton numbers L_e or L_μ, the current and interaction Lagrangian are given by

$$\begin{aligned} \mathcal{L}_{\text{int}} &= J^\mu_\Lambda\, \partial_\mu \phi / V_\Lambda + \text{quadratic terms in } \partial_\mu \phi\ , && (2)\\ J^\mu_\Lambda &= \Lambda_e(\bar{e}\,\gamma^\mu\, e + \bar{\nu}_e \gamma^\mu \nu_e) + \Lambda_\mu(\bar{\mu}\,\gamma^\mu\, \mu + \bar{\nu}_\mu \gamma^\mu \nu_\mu) + \cdots\ , && (3) \end{aligned}$$

where the dots stand for model dependent scalar boson contributions and higher order effective couplings.

All this can be easily understood [3, 4]. Let Λ_a and Λ_i be the quantum numbers of the fermion fields χ^a and scalars σ_i under an abelian global symmetry $\mathrm{U}(1)_\Lambda$ which is spontaneously broken by vacuum expectation values $\langle \sigma_i \rangle$. A suitable change of variables namely,

$$\begin{aligned} \chi^a &= e^{-i\,\Lambda_a \phi / V_\Lambda} \psi^a\ , && (4)\\ \sigma_i &= e^{-i\,\Lambda_i \phi / V_\Lambda}\ (\langle \sigma_i \rangle + \rho_i)\ , && (5) \end{aligned}$$

makes the Lagrangian to be expressed in terms of *physical weak* eigenstates, fermions $\psi^a = e,\ \nu_e,\ ...,$ massive bosons ρ_i and gauge bosons (assumed to be singlets of the global symmetry) all of them invariant under $\mathrm{U}(1)_\Lambda$. The only non-invariant field is the NG boson ϕ. Its couplings are derived from the kinetic Lagrangians of χ^a and σ_i by application of the transformations (4-5).

The second member of the NG equation of motion, Eq. (1), is clearly not an ordinary scalar density. If one calculates the divergency of a vector or axial-vector fermion current using the Dirac equation, one ends up with pseudo-scalars (both diagonal and non-diagonal in flavor) and *off-diagonal* scalar densities:

$$\partial_\mu \overline{f}_i \gamma^\mu \gamma_5 f_j = i(m_i + m_j)\overline{f}_i \gamma_5 f_j \ , \tag{6}$$

$$\partial_\mu \overline{f}_i \gamma^\mu f_j = i(m_i - m_j)\overline{f}_i f_j \ . \tag{7}$$

However couplings to *flavor-diagonal* scalar densities are not possible [3, 5]. The pseudo-scalars that are diagonal in flavor vanish for free particle states and the flavor violating densities depend on the relative phase of distinct flavors. The natural conclusion has been that these relative phases cancel each other when summed over a large number of particles and therefore long range '$1/r$' NG fields are not possible, only *spin-dependent* '$1/r^3$' potentials can exist [3, 5]. It will be clear in a moment that this is not true [4, 6, 7, 8].

2 Long range Nambu-Goldstone fields

First notice that the source of a NG field does not vanish if the matter current J^μ_Λ is not conserved. The volume integral

$$\int d^3x \ \ \partial_\mu J^\mu_\Lambda = \frac{d\Lambda^{\text{created}}}{dt} \tag{8}$$

represents the rate of *creation* per unity of time of the Λ-number carried by all particles except ϕ. This rate, divided by $-V_\Lambda$, constitutes the source of the NG boson and

$$\phi(t, \vec{r}) = \frac{-1}{V_\Lambda} \int d^3x \, \frac{\partial_\mu J^\mu_\Lambda(t - |\vec{r} - \vec{x}| \, , \vec{x})}{4\pi \, |\vec{r} - \vec{x}|} \tag{9}$$

is a solution of the equation of motion (1). It is a long-range field but suppressed by the ratio of the rate of Λ-violating processes per unity of time over the scale V_Λ. To produce a significant macroscopic field one needs thus a large rate of reactions. It calls for astrophysical sources. There is evidence that solar neutrinos change flavor as they come out of the Sun. If ν_e oscillates into ν_μ, for instance, then L_e and L_μ are not conserved in a large scale in stars and presumably in Supernovae as well. If any of these quantum numbers is associated with a spontaneously broken global symmetry then, the respective NG bosons (so-called majorons [5]) acquire classical long-range field configurations.

For definiteness suppose that $\Lambda = L_e - L_\mu$ is spontaneously broken and that a fraction of ν_es emitted from a star are resonantly converted into ν_μ through the Mikheyev-Smirnov-Wolfenstein (MSW) mechanism [9] (also $\nu_\mu \rightarrow \nu_e$ in a Supernova)

in a certain shell inside the star. For simplicity assume spherical symmetry and stationary neutrino fluxes in which case the $L_e - L_\mu$ number carried by electrons and neutrinos decreases at a constant rate in time. The rate $\dot{L}_e - \dot{L}_\mu$ generates a long-range NG field of Coulombian form, charge divided by $4\pi r$, above the resonance shell,

$$\phi = \frac{1}{4\pi r}\frac{2}{V_\Lambda}\left[\dot{N}(\nu_e \to \nu_\mu) - \dot{N}(\nu_\mu \to \nu_e)\right] \ , \tag{10}$$

whereas ϕ is constant below the resonance shell. This long-range '$1/r$' field configurations are in conflict with the conclusion traditionally taken by looking at the fermion bilinears (6-7) that appear as a source of a NG boson namely, that only '$1/r^3$' spin-dependent NG potentials could exist. This particular issue is the subject of a recent work [8] and is examined below.

In the interest of clearness let us concentrate on a specific system, the $\nu_e - \nu_\mu$ pair of flavors. Their interactions are assumed to be weak interactions with a background medium, couplings with a NG boson associated to $\Lambda = L_e - L_\mu$ and a 2×2 *real* Majorana mass matrix m, all contained in:

$$\begin{aligned}\mathcal{L}_{\rm int} \ &= \ -\sqrt{2}G_{\rm F} n_e(\bar{\nu}_e\gamma^0\nu_e) - \frac{g\langle Z^0\rangle}{2\cos\theta_{\rm W}}(\bar{\nu}_e\gamma^0\nu_e + \bar{\nu}_\mu\gamma^0\nu_\mu) \\ &+ \ \frac{1}{V_\Lambda}\partial_\mu\phi(\bar{\nu}_e\gamma^\mu\nu_e - \bar{\nu}_\mu\gamma^\mu\nu_\mu) - (\nu_L^T C m \nu_L + \text{H. C.}) \ .\end{aligned} \tag{11}$$

n_e is the electron density. Notice that only m_{ee} and $m_{\mu\mu}$ break $L_e - L_\mu$. This Lagrangian determines the equations of motion of the (left-handed) neutrino operators and also of the single particle wave functions, related to them by a relation of the type $\nu_L + \nu_L^C = \int a\,\psi + a^\dagger\,\psi^C$. The wave functions obey the equations

$$i\,\partial\!\!\!/\psi_a = m_{ab}\,\psi_b - V_a^\mu\,\gamma_\mu\gamma_5\psi_a \ , \tag{12}$$

where a stands for ν_e and ν_μ and V_a^μ represent the *flavor conserving* weak and majoron vector potentials ($V_a^\mu = -\Lambda_a\partial^\mu\phi/V_\Lambda$ in the latter case). When $\partial_\mu J_\Lambda^\mu$ is evaluated over a system of neutrino particles the equation of motion (1) becomes [8]:

$$\partial_\mu\partial^\mu\phi = \frac{2}{V_\Lambda}\sum_\nu i\left(m_{ee}\,\bar{\psi}_{\nu_e}\gamma_5\psi_{\nu_e} - m_{\mu\mu}\,\bar{\psi}_{\nu_\mu}\gamma_5\psi_{\nu_\mu}\right) \ . \tag{13}$$

As expected, the parameters m_{ee} and $m_{\mu\mu}$ that break $L_e - L_\mu$ in the Lagrangian (11) appear explicitly in the source terms. These are precisely the kind of pseudo-scalar densities identified in Eq. (6). They would vanish identically if the neutrinos were free particles. But they are not. When studying neutrino oscillations one usually separates the spin from the flavor degrees of freedom by expressing the wave function as a product of a left-handed *free massless* spinor ψ_0^α ($\gamma_5\psi_0 = -\psi_0$, $i\,\partial\!\!\!/\psi_0 = 0$) and

a flavor-valued wave function φ_a. It obeys a well known evolution equation [10] as a function of the distance s travelled by each particle,

$$i\frac{d}{ds}\left(\begin{array}{c}\varphi_e\\ \\ \varphi_\mu\end{array}\right)=\frac{1}{2E}\left(\begin{array}{cc}m^2_{ee}+2EV_e & \frac{1}{2}\Delta m^2\sin 2\theta\\ \\ \frac{1}{2}\Delta m^2\sin 2\theta & m^2_{\mu\mu}+2EV_\mu\end{array}\right)\left(\begin{array}{c}\varphi_e\\ \\ \varphi_\mu\end{array}\right)\ , \tag{14}$$

see Eqs. (18-21). However, pure chiral spinors give vanishing scalar densities in Eq. (13) so one has to go beyond the approximation $\psi_a=\psi_0\varphi_a$.

For a spherical symmetric NG field as in Eq. (10) the velocity of a radially moving neutrino is parallel to the vector potentials $\mathbf{V}_a\propto\nabla\phi$ and in that case an approximate solution of Eq. (12) is [8]

$$\psi^\alpha_a(x)\cong\psi^\alpha_0\,\varphi_a+\gamma^0_{\alpha\beta}\frac{m_{ab}}{2E}\,\psi^\beta_0\,\varphi_b\ . \tag{15}$$

Applying to the majoron equation of motion (13) it becomes

$$\begin{aligned}\partial_\mu\partial^\mu\phi &= \frac{1}{V_\Lambda}\sum_\nu\frac{m_{e\mu}}{E}(m_{ee}+m_{\mu\mu})\,i\,(\varphi^\dagger_e\varphi_\mu-\varphi^\dagger_\mu\varphi_e)\,\psi^\dagger_0\psi_0\\ &= \frac{1}{V_\Lambda}\sum_\nu\frac{\Delta m^2}{2E}\sin 2\theta\,i\,(\varphi^\dagger_1\varphi_2-\varphi^\dagger_2\varphi_1)\,\psi^\dagger_0\psi_0\ .\end{aligned} \tag{16}$$

where $\varphi_{1,2}$ refer to the mass eigenstate basis. It is now clear that the generation of a NG field requires the interference between different flavor components of the ν wave functions. But that simply means the existence of neutrino oscillations. A crucial fact is that the source term does not depend on arbitrary initial phases of the ν wave functions. Hence, there is no automatic phase cancellation when the sum is taken over a large number of particles. This of course makes possible to have sources of macroscopic dimensions.

Making use of Eq. (14), the last result can be written in terms of the probabilities of observing the ν_e and ν_μ flavors respectively, $P_{\nu_e}=\varphi^\dagger_e\varphi_e/\varphi^\dagger\varphi$ and $P_{\nu_\mu}=1-P_{\nu_e}$:

$$\partial_\mu\partial^\mu\phi_\Lambda=\frac{-1}{V_\Lambda}\sum_\nu\psi^\dagger\psi\,\frac{d(P_{\nu_e}-P_{\nu_\mu})}{ds}=\frac{-1}{V_\Lambda}\frac{d(L_e-L_\mu)}{dt\,d^3x}\ . \tag{17}$$

Again, the second member does not depend on the initial phases of the individual particles. It also tells under which conditions a large number of particles makes a charge with definite sign: it is when the $\Lambda=L_e-L_\mu$ number carried by neutrinos suffers a net increase or decrease in Λ-violating processes. We come back to the point of view at the beginning of this section: the source term of a NG field is proportional to the rate of *creation* (not just local variation) per unity of time and volume of the quantum number associated to it that is carried by the matter particles [4, 6, 7]. In a star or reactor where only ν_e are produced out of electrons the oscillations $\nu_e\to\nu_\mu$ necessarily lead to a total decrease of the lepton number L_e-L_μ: the electrons captured in nuclear reactions are, after all, totally or partially transformed into ν_μ (the probability of observing the ν_μ flavor is not zero).

3 Neutrino oscillations: role of NG fields

Understood as they are the conditions under which a long-range NG field is produced the next question is what are the implications? A NG field only interacts through its gradient which implies a suppression factor of V_Λ^{-1} over the time-length scale of field variation. This is typically an extremely small number. However ν oscillations are sensitive to very tiny external potentials that are not independent of the flavor. This condition is satisfied by the majoron couplings as a rule, see Eqs. (2-3), for example. They contribute with the vector potentials $V_a^\mu = -\Lambda_a \partial^\mu \phi / V_\Lambda$ to the fermion equations of motion (12) and potentials

$$V_a = -\Lambda_a (\dot{\phi} + \mathbf{v} \cdot \nabla \phi)/V_\Lambda \; , \tag{18}$$

where $\mathbf{v}$ is the neutrino velocity, to the flavor oscillation equation (14).

In the case analyzed in the previous section, the important quantity for $\nu_e \leftrightarrow \nu_\mu$ oscillations is the difference

$$V_e - V_\mu = \sqrt{2} G_{\rm F} n_e + \frac{4}{V_\Lambda^2} \frac{1}{4\pi r^2} \left[\dot{N}(\nu_e \to \nu_\mu) - \dot{N}(\nu_\mu \to \nu_e) \right] \; . \tag{19}$$

It indicates what is the typical order of magnitude of the NG potentials ($V_{\rm NG}$), neutrino flux over V_Λ^2. In the case of the Sun, $j_\nu / V_\Lambda^2 \sim 10^{-2}/R_\odot$ at the Sun radius for a scale V_Λ as low as 1 KeV. It looks too small to perturb solar neutrino oscillations. In a Supernova however the fluxes are many orders of magnitude higher and

$$V_{\rm NG} \sim \frac{j_\nu}{V_\Lambda^2} = 1.48 \, \frac{V_\Lambda^{-2}}{G_{\rm F}} \frac{L_\nu}{10^{52}\,\text{ergs/s}} \frac{10\,\text{MeV}}{\langle E_\nu \rangle} \left(\frac{r}{10^{10}\,\text{cm}} \right)^{-2} \times 10^{-12} \text{ eV} \; , \tag{20}$$

where L_ν is the ν energy luminosity, $\langle E_\nu \rangle$ the average energy and $L_\nu / \langle E_\nu \rangle$ the neutrino emission rate. In turn, the electroweak potential [9] is given in a Supernova envelope star by

$$V_W = \sqrt{2}\, G_{\rm F}\, n_e = 0.76 \; Y_e \frac{\tilde{M}}{10^{31}\text{g}} \left(\frac{r}{10^{10}\,\text{cm}} \right)^{-3} \times 10^{-12} \text{ eV} \; , \tag{21}$$

where $\tilde{M} = \rho\, r^3$ is a constant between 1 and 15 $\times 10^{31}$g, depending on the star [11].

Clearly, the NG potentials compete with the electroweak potentials for scales V_Λ as high as $G_{\rm F}^{-1/2}$ and due to their long-range nature even overcome at large enough radius. One can conceive that a NG field may affect the propagation of other neutrinos or flavors outside the region it was created or even outside the star. On the other hand, it may participate in resonant oscillations for values of Δm^2 that are interesting for solar or atmospheric neutrino solutions. It was shown in two previous papers that NG fields may improve the adiabaticity of $\nu_e \leftrightarrow \nu_X$ transitions [6] or cause *resonant* anti-neutrino oscillations otherwise impossible with solely weak interactions [4].

3.1 Improvement of adiabaticity: majoron back reaction

Suppose that ν_e oscillates in the Sun into an active neutrino in accordance with the non-adiabatic, small mixing angle solution [12, 13]. Then, the same $\nu_e \to \nu_\mu$ transitions must occur in a Supernova envelope, non-adiabatic as well, but accompanied by $\nu_\mu \to \nu_e$ since ν_μ are also produced in a Supernova (for convenience ν_μ designates the relevant active neutrino). First, as results from the resonance condition,

$$2E(V_e - V_\mu) = \Delta m^2 \cos 2\theta \ , \tag{22}$$

the neutrinos with higher energy are converted at a larger radius. Second, the transitions are less adiabatic and less probable for the higher energy neutrinos [14] which has important consequences: since the $\nu_{\mu,\tau}$ flavors are emitted with an hotter spectrum than ν_e, the final observable ν_e spectrum is softer than if the transitions were completely adiabatic.

The situation changes if there is a majoron field associated to a partial lepton number, let it be L_e [6] or $L_e - L_\mu$ as considered in the previous sections. The potential $V_e - V_\mu$ departs from the charged current potential as shows Eq. (19). The rates $\dot{N}(\nu_e \to \nu_\mu)$ and $\dot{N}(\nu_\mu \to \nu_e)$ are functions of the radius, in other words, the ν conversions inside a certain sphere create a NG field gradient outside that sphere. Since a Supernova emits more ν_e than ν_μ the majoron potential is positive which has the effect of slowing down the fall of $V_e - V_\mu$ with the radius and, by virtue of Eq. (22), of increasing the radius at which the higher energy neutrinos are converted. For a fixed ν energy, the non-adiabaticity increases with the slope $|d\ln(V_e - V_\mu)/dr|$ [10, 12] thus, the net effect of the majoron field is to improve the adiabaticity and conversion efficiency of the hottest neutrinos. The implication, observable in detectors like Super-Kamiokande and SNO [15], is a ν_e spectrum harder than expected within the framework of electroweak interactions and non-adiabatic solar neutrino solution.

3.2 Resonant conversion of anti-neutrinos

If not only one but all three partial lepton numbers are spontaneously broken one expects to exist mixing between the three NG bosons ξ_e, ξ_μ, ξ_τ [4]. By going to the unitary basis as in Eqs. (4-5),

$$\chi^a = \exp\{-i(\xi_e L_e^a + \xi_\mu L_\mu^a + \xi_\tau L_\tau^a)\}\, \psi^a \ , \tag{23}$$

$$\sigma_i = \exp\{-i(\xi_e L_e^i + \xi_\mu L_\mu^i + \xi_\tau L_\tau^i)\}\, (\langle\sigma_i\rangle + \rho_i) \ , \tag{24}$$

the χ^a fermion and σ_i boson kinetic terms deliver the following interaction Lagrangian

$$\mathcal{L}_{\rm int} = \frac{1}{2} V_{\ell m}^2\, \partial_\mu \xi_\ell\, \partial^\mu \xi_m + \partial_\mu \xi_\ell (\bar{\ell}\, \gamma^\mu\, \ell + \bar{\nu}_\ell \gamma^\mu \nu_\ell) \ , \qquad \ell, m = e, \mu, \tau \ . \tag{25}$$

The kinetic matrix $V^2_{\ell m}$ and its inverse $G_{\ell m}$ (same dimension as $G_{\rm F}$) are not diagonal in general [4]. As a result, long-range majoron fields couple different flavors to each other. Instead of an equation like (19), a neutrino with flavor ν_ℓ feels a NG potential

$$V_\ell = -(G_{\ell e}\,\dot{L}_e + G_{\ell\mu}\,\dot{L}_\mu + G_{\ell\tau}\,\dot{L}_\tau)/4\pi r^2\ . \tag{26}$$

It means that the violation of one flavor in one place creates majoron fields that act at a distance on any of the other flavors. This may have interesting consequences.

Suppose that resonant $\nu_e \leftrightarrow \nu_H$ oscillations take place in a Supernova environment (H stands for Heavy, it may be the heaviest of the ν mass eigenstates involved in atmospheric ν oscillations.) In addition ν_e also mixes with ν_L (L means Light, possibly the neutrino involved in solar ν oscillations) and, in absence of other than weak interactions, ν_e oscillates resonantly into ν_L, but not $\bar{\nu}_e \to \bar{\nu}_L$, in a region of lower density than where $\nu_e \leftrightarrow \nu_H$ take place. If majoron fields come into play, the ν_e and ν_L propagation is modified by the potentials in Eq. (26) that are created by the $\nu_e \leftrightarrow \nu_H$ oscillations, with $\dot{L}_e = -\dot{L}_\mu < 0$ because a Supernova emits more ν_es than ν_Hs. If the following hierarchy, $|G_{e\ell}| << |G_{LH}|, G_{HH}$, exists between these constants, the relevant potential for $\bar{\nu}_e \leftrightarrow \bar{\nu}_L$ oscillations is approximately

$$V_{\bar{\nu}_e} - V_{\bar{\nu}_L} = -\sqrt{2}\,G_{\rm F}\,n_e\ - G_{LH}\,\dot{L}_H/4\pi r^2 = -(V_{\nu_e} - V_{\nu_L})\ . \tag{27}$$

In absence of NG fields this potential is negative and, by hypothesis, the resonance condition (22) cannot be satisfied for $\bar{\nu}_e \leftrightarrow \bar{\nu}_L$ contrary to $\nu_e \leftrightarrow \nu_L$. However, if the majoron fields exist and G_{LH} is negative, the NG potential is positive and necessarily overcomes at large enough radius because n_e falls faster than $1/r^2$. Therefore, if Δm^2_{eL} is sufficiently low and the ν fluxes high enough resonant anti-neutrino oscillations become possible. Such effect can in principle be observed because the predicted energy spectrum of $\bar{\nu}_\mu$ and $\bar{\nu}_\tau$ is harder than the $\bar{\nu}_e$ spectrum.

The SN1987A data [16] seem to disfavor $\bar{\nu}_e \leftrightarrow \bar{\nu}_{\mu,\tau}$ oscillations [17] (see however [18]) but do not exclude that they occur with a probability below a certain limit. As stressed in [4], they may occur in the first instants of ν emission, when the ν luminosity and hypothetical majoron fields are stronger, to disappear later below some ν flux threshold. This kind of time dependence, if observed, in correlation with the magnitude of ν fluxes, would be a clear signature of majoron fields providing in addition a good measurement of the lepton symmetry breaking.

4 Conclusions

In spite of the fact that true Nambu-Goldstone bosons only have derivative couplings it is still possible to obtain macroscopic long-range NG fields. The source of a NG field is the time rate of decreasing of the quantum number associated with it that is carried by the matter particles. If the lepton numbers are spontaneously

broken, neutrino oscillations give rise to long-range majoron fields as a result of the constructive interference between the wave functions of different flavors. If the scale of lepton symmetry breaking is at or below the Fermi scale the associated majoron fields are significant enough to change the ν oscillations in Supernovae (SN). That corresponds to a sensitivity to effective neutrino-majoron coupling constants in scattering processes as low as $m_\nu G_{\mathrm{F}}^{1/2}$! One of the imprints of majoron fields is the *surprise*: because in solar, atmospheric and terrestrial neutrino experiments the fluxes are probably too low (or too high the scale of lepton symmetry breaking) to produce significant majoron fields, their effects in SN neutrinos cannot be anticipated from those kind of experiments. The other point is the correlation with the magnitude of ν fluxes. It tells what is the scale of lepton symmetry breaking, and in a more refined way if, as a result of the few seconds decay of SN ν emission, the majoron fields and observed SN ν oscillation patterns exhibit a time dependence in correlation with the ν luminosity.

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